CMU-HEP95-11
DOE-ER/40682-101
13A-40561-INT95-017
DOE Research and
Development Report

Charmed Baryon Masses in Chiral Perturbation Theory *

Martin J. Savage†
Department of Physics, Carnegie Mellon University
Pittsburgh, Pennsylvania 15213 U.S.A.
savage@thepub.phys.cmu.edu
(August 1995)

Abstract

The masses of the charmed baryons in the **6** representation of $SU(3)$ obey an equal spacing rule at lowest order in $SU(3)$ breaking, $\mathcal{O}(m_s)$. We compute the corrections to this relation at order $\mathcal{O}(m_s^{3/2})$ arising from meson loops using chiral perturbation theory combined with heavy quark symmetry and find them to be small. We also examine the hyperfine interaction responsible for the splitting between the $J^\pi = \frac{3}{2}^+$ and $J^\pi = \frac{1}{2}^+$ baryons in the **6** representation. The results also hold in the b-baryon sector.

*Work supported in part by the U.S. Dept. of Energy under Grant No. DE-FG02-91-ER40682

†DOE Outstanding Junior Investigator

Our knowledge of the masses and properties of the lowest lying charmed baryons has improved dramatically during the last few years [1]. With the recent discovery of the Ξ_{c2}^{0*} ($J^\pi = \frac{3}{2}^+$) of mass 2642.8 ± 2.2 MeV [2] and hints of the Ξ_{c2}^+ ($J^\pi = \frac{1}{2}^+$) with a mass of ~ 2560 MeV [3] and Σ_c^{++*} ($J^\pi = \frac{3}{2}^+$) with a mass of ~ 2530 MeV [4] it has become timely to see just how well we understand the pattern of masses in the charmed baryon sector. There have been many estimates of the charmed baryon masses made in the past [5]. Recently, Rosner [6] has performed a spin-flavour analysis of the charmed baryons and the lowest lying noncharmed baryons to obtain masses and strong decay widths. He obtained mass relations between the $J^\pi = \frac{3}{2}^+$ and $J^\pi = \frac{1}{2}^+$ charmed baryons in the lowest lying $\mathbf{6}$ representations including SU(3) breaking arising from the difference between the strange and non-strange constituent quark masses. In this work we will examine SU(3) breaking in the lowest lying $\mathbf{6}$ representation of charmed baryons using chiral perturbation theory with heavy quark symmetry. At lowest order in SU(3) breaking there is an equal-spacing rule analogous to the equal spacing rule in the noncharmed baryon decuplet that receives finite and computable corrections. We determine that these corrections are small. The hyperfine mass splittings between the $\mathbf{6}$ and $\mathbf{6}^*$ are also examined.

Heavy quark symmetry and chiral symmetry are combined together in order to describe the soft hadronic interactions of hadrons containing a heavy quark [7,8]. The light degrees of freedom in the ground state of a baryon containing one heavy quark can have $s_l = 0$ corresponding to a member of the flavour $SU(3)$ $\overline{\mathbf{3}}$, $T_i(v)$, with $J^\pi = \frac{1}{2}^+$ or they can have $s_l = 1$ corresponding to a member of the flavour $SU(3)$ $\mathbf{6}$, $S_\mu^{ij}(v)$. In the latter case, the spin of the light degrees of freedom can be combined with the spin of the heavy quark to form both $J^\pi = \frac{3}{2}^+$ and $J^\pi = \frac{1}{2}^+$ baryons, which are degenerate in the $m_Q \to \infty$ limit. Baryons in the $\overline{\mathbf{3}}$ and $\mathbf{6}$ representations are described by the fields

$$\begin{aligned} S_\mu^{ij}(v) &= \tfrac{1}{\sqrt{3}}(\gamma_\mu + v_\mu)\gamma_5 \tfrac{1}{2}(1+\not{v})B^{ij} \ + \ \tfrac{1}{2}(1+\not{v})B_\mu^{*ij} \\ T_i(v) &= \tfrac{1}{2}(1+\not{v})B_i \quad , \end{aligned} \tag{1}$$

where the $J^\pi = \frac{1}{2}^+$ charmed baryons of the $\mathbf{6}$ are assigned to the symmetric tensor B^{ij}

$$\begin{aligned} B^{11} &= \Sigma_c^{++} \ , \ B^{12} = \tfrac{1}{\sqrt{2}}\Sigma_c^+ \ , \ B^{22} = \Sigma_c^0 \ , \\ B^{13} &= \tfrac{1}{\sqrt{2}}\Xi_{c2}^+ \ , \ B^{23} = \tfrac{1}{\sqrt{2}}\Xi_{c2}^0 \ , \ B^{33} = \Omega_c^0 \quad . \end{aligned} \tag{2}$$

The $J^\pi = \frac{3}{2}^+$ partners of these baryons have the same $SU(3)$ assignment in B_μ^{*ij}. The charmed baryons of the $\overline{\mathbf{3}}$ representation are assigned to B_i as

$$B_1 = \Xi_{c1}^0 \ , \ B_2 = -\Xi_{c1}^+ \ , \ B_3 = \Lambda_c^+ \quad . \tag{3}$$

The chiral lagrangian describing SU(3) invariant soft hadronic interactions of the charmed baryons is [8]

$$\begin{aligned} \mathcal{L}_Q = {} & i\overline{T}^i v\cdot DT_i \ - \ i\overline{S}_{ij}^\mu v\cdot DS_\mu^{ij} \ + \ \Delta_0 \overline{T}^i T_i + \tfrac{f^2}{8} Tr\left[\partial^\mu\Sigma\partial_\mu\Sigma^\dagger\right] \\ & + \ g_3\left(\epsilon_{ijk}\overline{T}^i (A^\mu)^j_l S_\mu^{kl} + h.c.\right) + \ ig_2\ \epsilon_{\mu\nu\rho\sigma}\overline{S}_{ik}^\mu v^\nu (A^\rho)^i_j S^{\sigma jk} \ + \ \cdots \quad , \end{aligned} \tag{4}$$

where the dots denote operators with more derivatives or those that are higher order in the $1/m_Q$ expansion and D^α is the chiral covariant derivative. The axial chiral field

$A^\mu = \frac{i}{2}\left(\xi^\dagger\partial^\mu\xi - \xi\partial^\mu\xi^\dagger\right)$ is defined in terms of $\xi = \exp(iM/f)$ where M is the octet of pseudo-Goldstone bosons

$$M = \begin{pmatrix} \frac{1}{\sqrt{6}}\eta + \frac{1}{\sqrt{2}}\pi^0 & \pi^+ & K^+ \\ \pi^- & \frac{1}{\sqrt{6}}\eta - \frac{1}{\sqrt{2}}\pi^0 & K^0 \\ K^- & \overline{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix} \quad , \tag{5}$$

and $f \sim 132$MeV is the pion decay constant at lowest order. The Σ field of pseudo-Goldstone bosons is $\Sigma = \xi^2 = \exp(i2M/f)$. Coupling of a single pseudo-Goldstone boson to the $\overline{\mathbf{3}}$ baryons is forbidden at lowest order in $1/m_Q$. Even in the infinite mass limit the $\mathbf{6}$ baryons are not degenerate with the $\overline{\mathbf{3}}$ baryons as the light degrees of freedom are in a different configuration giving rise to an intrinsic mass difference Δ_0. We have chosen to remove the mass of the $\mathbf{6}$ from the fields and not the mass of the $\overline{\mathbf{3}}$ for convenience. The masses of the charmed baryons that follow from Eq. (4) are trivial in the sense that there is no SU(3) breaking and the charmed baryons in the $\overline{\mathbf{3}}$ have equal mass, as do those in the $\mathbf{6}$ but the $\overline{\mathbf{3}}$ and $\mathbf{6}$ are split by Δ_0.

The strong coupling constants g_2 and g_3 must be determined from experimental data on the strong widths or from loop processes. Observation of the Ξ_{c2}^{0*} and the upper limit on its width [2] of $\Gamma(\Xi_{c2}^{0*}) < 5.5$ MeV constrains g_3 (neglecting higher order corrections) to be $|g_3| < 1.3$. The coupling constant g_2 is, as yet, unconstrained. We notice that the upper bound on g_3 is already below the value one would expect from large N_c considerations [9,10], $g_2 = -\frac{3}{2}g_A = -1.9$ and $g_3 = \sqrt{\frac{3}{2}}g_A = 1.5$ with $g_A \sim 1.25$.

SU(3) breaking in the masses of the charmed baryons arises from explicit insertions of the light quark mass matrix and from loop graphs involving the pseudoscalar mesons. The general form of such corrections is discussed in [11] ; however, we wish to be more specific. Using the notation of [11] we write the lagrange density that is linear in the light quark mass matrix and is lowest order in the heavy quark expansion

$$\mathcal{L} = \lambda_1 \overline{S}^\mu_{ij}(\chi^+)^i_k S^{jk}_\mu + \lambda_2 \overline{S}^\mu_{ij} S^{ij}_\mu (\chi^+)^k_k + \lambda_3 \overline{T^i}(\chi^+)^j_i T_j + \lambda_4 \overline{T^i} T_i (\chi^+)^k_k - \mu(\chi^+)^k_k \quad , \tag{6}$$

where

$$\chi^+ = \xi^\dagger m_q \xi^\dagger + \xi m_q \xi \quad , \tag{7}$$

and where the light quark mass matrix is

$$m_q = \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_d & 0 \\ 0 & 0 & m_s \end{pmatrix} \quad . \tag{8}$$

The last term in Eq. (6) generates the non-zero masses of the pseudo-Goldstone bosons while the first four terms contribute to the masses of the charmed baryons. Non-zero values for λ_1 and λ_3 give rise to the leading SU(3) breaking between charmed baryon masses. Performing the contraction of indices in Eq. (6) one generates baryon masses in the presence of octet $SU(3)$ breaking. In general, the $\mathbf{6}$ could have breaking terms with representations $\mathbf{6}\otimes\overline{\mathbf{6}} = \mathbf{27}\oplus\mathbf{8}\oplus\mathbf{1}$ but the single insertion of m_q gives octet breaking only. As there are only three elements in the $\overline{\mathbf{3}}$ and two are degenerate in the limit of isospin symmetry there are no $SU(3)$ mass relations that hold in the presence of Eq. (6). However, there is a nontrivial mass relation between baryons in the **6** that holds in the presence of Eq. (6),

$$\frac{1}{3}\left(M_{\Sigma_c^{++}}+M_{\Sigma_c^{+}}+M_{\Sigma_c^{0}}\right)+M_{\Omega_c^0}-\left(M_{\Xi_{c2}^{+}}+M_{\Xi_{c2}^{0}}\right)=0 \quad , \tag{9}$$

which in the limit of isospin symmetry becomes

$$M_{\Sigma_c}+M_{\Omega_c}-2M_{\Xi_{c2}}=0 \quad . \tag{10}$$

This is an equal spacing rule analogous to the equal spacing rule that arises in the decuplet of uncharmed $J^\pi=\frac{3}{2}^+$ baryons. We have not yet discussed mixing between the Ξ_c's in the $\overline{\mathbf{3}}$ and **6** . Mixing between these particles is both SU(3) breaking and a $1/M_Q$ effect (as it requires mixing between states with $s_l=0$ and $s_l=1$ in the heavy quark limit). Further, the mixing term will enter squared when the $\{\Xi_{c1},\Xi_{c2}\}$ mass matrix is diagonalized and therefore we neglect it.

Corrections to the equal spacing rule Eq. (10) arise from more insertions of the light quark mass matrix and from loops involving the pseudo-Goldstone bosons. The leading corrections arise from the meson loops and are of order $m_s^{3/2}$. Graphs involving the $\mathbf{6}^{(*)}$ baryons depend on the meson masses only (neglecting the $\mathcal{O}(m_s)$ splittings between the intermediate state baryons within the loop that give corrections higher order in m_s), while those involving the $\overline{\mathbf{3}}$ baryons are functions of the meson masses and the mass splitting Δ_0. Despite the loop contribution to the individual masses being of order a few hundred MeV (expressions for which can be derived from results in [11]), the correction to the equal spacing rule is small; explicitly we find that

$$M_{\Sigma_c}+M_{\Omega_c}-2M_{\Xi_c}=\frac{1}{48\pi f^2}\left[g_2^2\mathcal{J}(0)-g_3^2\mathcal{J}(\Delta_0)\right] \quad , \tag{11}$$

where the function $\mathcal{J}(y)$ is given by

$$\begin{aligned}\mathcal{J}(y)=\tfrac{1}{\pi}\Big[&-y^3\log\left(m_K^8/m_\eta^6 m_\pi^2\right)+6ym_K^2\log\left(m_K^2/m_\pi^2\right)-\tfrac{9}{2}ym_\eta^2\log\left(m_\eta^2/m_\pi^2\right)\\ &-4\mathcal{G}(y,m_K)+3\mathcal{G}(y,m_\eta)+\mathcal{G}(y,m_\pi)\Big]\end{aligned} \tag{12}$$

and

$$\mathcal{G}(y,m)=(y^2-m^2)^{\frac{3}{2}}\log\left(\frac{-y-\sqrt{y^2-m^2+i\epsilon}}{-y+\sqrt{y^2-m^2+i\epsilon}}\right) \quad . \tag{13}$$

In order to arrive at this result we have used the Gell-Mann–Okubo mass relation between the pseudo-Goldstone boson masses $4m_k^2-3m_\eta^2-m_\pi^2=0$ that arises at lowest order from Eq. (4) and Eq. (6). Notice that the correction to the equal-spacing rule Eq. (10) (and Eq. (9)) that holds for $\mathbf{8}\oplus\mathbf{1}$ $SU(3)$ breaking is finite. This results from the fact that any corrections to Eq. (10) and Eq. (9) must transform as a **27** under $SU(3)$ and there are no counterterms in Eq. (4) and Eq. (6) that transform as a **27** to absorb divergences (such counterterms start at $\mathcal{O}(m_s^2)$). In the limit of vanishing **6** - $\overline{\mathbf{3}}$ mass splitting Δ_0 we have

$$\mathcal{J}(0)=4m_K^3-3m_\eta^3-m_\pi^3 \quad . \tag{14}$$

This particular combination of masses appears in the violation of mass relations that hold in the presence of octet $SU(3)$ breaking in other hadronic sectors, in the octet baryon sector [12] and in the vector meson sector [13]. This is because a **27** representation is required to violate the mass relations in each sector and the combination of masses that appears in Eq. (14) is the unique combination that transforms as a **27** [13,14]. Numerically, we find that the right hand side of Eq. (11) is very small $\sim 5(g_2^2 - g_3^2)$MeV using the masses of the charged K and π and setting $\Delta_0 = 100$ MeV (the result is very insensitive to the value of Δ_0). Therefore, we expect that the equal-spacing rule in Eq. (10) is well satisfied. We can use the mass of the Σ_c^{++} , 2453.1 ± 0.6 MeV and the mass of the Ω_c^0, 2704 ± 4 MeV to predict that

$$\begin{aligned} M_{\Xi_{c2}} &= \tfrac{1}{2}\left[M_{\Sigma_c} + M_{\Omega_c}\right] \\ &\sim 2579 \text{ MeV} \quad , \end{aligned} \tag{15}$$

which we expect to be within a few MeV of the actual mass. This is in contrast to the recent experimental suggestion [3] of $M_{\Xi_{c2}} \sim 2560$ MeV, some 20 MeV away from Eq. (15). The mixing between the **6** and $\overline{\mathbf{3}}$ that we have neglected in our analysis will only increase the mass computed in Eq. (15), further increasing the possible discrepancy. It seems best not to consider this a serious problem or to compute higher order corrections to Eq. (10) until the experimental situation becomes more certain.

Turning now to the hyperfine mass splitting between the **6** and $\mathbf{6}^*$. Such a splitting results from the charm quark not being infinitely more massive than the scale of strong interactions. One can make a crude estimate for the magnitude of the splitting of $\delta_6 \sim \Lambda_{\rm QCD}^2/m_c \sim$ 50 MeV. Using the Ξ_{c2}^{0*} mass measurement [2] and the Ξ_{c2} mass determined from the Eq. (15) we find $\delta_6 \sim$ 64 MeV, consistent with our naive estimate. The hyperfine mass splittings are induced at lowest order by the $SU(3)$ invariant lagrange density [15]

$$\mathcal{L} = \frac{\delta_6}{6}(g_{\mu\alpha}g_{\nu\beta} - g_{\nu\alpha}g_{\mu\beta})\overline{S}^{\mu}_{ij} i\sigma^{\alpha\beta} S^{\nu ij} \quad . \tag{16}$$

It is clear that this operator gives rise to an equal-spacing rule for the hyperfine mass splittings (we use $\delta_{\Sigma_c} = M_{\Sigma_c^*} - M_{\Sigma_c}$, $\delta_{\Xi_c} = M_{\Xi_c^*} - M_{\Xi_c}$ and $\delta_{\Omega_c} = M_{\Omega_c^*} - M_{\Omega_c}$ for compactness)

$$\delta_{\Sigma_c} = \delta_{\Xi_c} = \delta_{\Omega_c} \quad , \tag{17}$$

from which we predict that

$$M_{\Sigma_c^*} \sim 2518 \text{ MeV} \quad \text{and} \quad M_{\Omega_c^*} \sim 2768 \text{ MeV} \quad . \tag{18}$$

The values in Eq. (18) are both within 4 MeV of the masses predicted by Rosner [6] using spin-flavour wavefunctions. Note that we have used the hyperfine mass splitting based on the equal-spacing rule prediction for the mass of the Ξ_{c2}. If we had used the value suggested by [3] then the predicted masses of the Σ_c^* and Ω_c^* would be ~ 2538 MeV and ~ 2788 MeV respectively.

When we consider $SU(3)$ breaking of the hyperfine mass splittings, an equal spacing rule analogous to that in Eq. (10) follows at linear order in m_s (assuming isospin symmetry),

$$\delta_{\Sigma_c} + \delta_{\Omega_c} - 2\delta_{\Xi_c} = 0 \quad , \tag{19}$$

from the lagrange density

$$\mathcal{L} = \frac{\delta_6'}{6}(g_{\mu\alpha}g_{\nu\beta} - g_{\nu\alpha}g_{\mu\beta})\overline{S}^{\mu}_{ij} i\sigma^{\alpha\beta} S^{\nu ik}(\chi^+)^j_k + \frac{\delta_6''}{6}(g_{\mu\alpha}g_{\nu\beta} - g_{\nu\alpha}g_{\mu\beta})\overline{S}^{\mu}_{ij} i\sigma^{\alpha\beta} S^{\nu ij}(\chi^+)^k_k \quad . \tag{20}$$

As the hyperfine mass splittings are a $1/m_c$ effect the leading $SU(3)$ breaking corrections to these splittings from meson loops must vanish as δ_6 vanishes. This means that the leading corrections to Eq. (19) are not $m_s^{3/2}$ but $m_s \log m_s$, when δ_6 is treated as small. Explicit computation of the pseudo-Goldstone boson loop graphs (setting $\Delta_0 = 0$) gives a finite correction to Eq. (19) of

$$\delta_{\Sigma_c} + \delta_{\Omega_c} - 2\delta_{\Xi_c} = \frac{\delta_6(3g_2^2 - 2g_3^2)}{16\pi^2 f^2}\left[m_k^2 \log\left(\frac{m_K^2}{m_\pi^2}\right) - \frac{3}{4}m_\eta^2 \log\left(\frac{m_\eta^2}{m_\pi^2}\right)\right] \quad , \tag{21}$$

where we have again used the Gell-Mann–Okubo mass formula for the mesons to simplify the expression. We see that the correction to the hyperfine equal-spacing rule in Eq. (21) is small ($\sim 10^{-2}\delta_6(3g_2^2 - 2g_3^2)$) and therefore we expect the equal-spacing rule to be well satisfied.

Unlike the $SU(3)$ corrections to the individual baryon masses from meson loops in the heavy quark limit which are finite when $\Delta_0 \to 0$, the loop corrections to the individual hyperfine splittings are divergent, and require presently unknown counterterms to absorb the divergence. The nonanalytic contributions from the loop graphs are of the form $m_M^2 \log(m_M/\Lambda_\chi)$ where we have renormalized at the chiral symmetry breaking scale Λ_χ. The same type of corrections contribute to the hyperfine mass splittings between the vector and pseudoscalar mesons containing a heavy quark. It is found that the counterterm required to reproduce the observed spectrum essentially exactly cancels the contribution of the chiral logarithm [16]. This leads one to believe that chiral perturbation theory, in particular the neglect of the counterterms, may be failing for the hyperfine mass splittings. However, as the hyperfine equal-spacing rule in Eq. (21) is independent of the counterterms, one might hope that the loop graphs do give a reasonable estimate of the size of the violation.

In conclusion, there is an equal-spacing rule that holds for the masses of the charmed baryons in the **6** representation of $SU(3)$ in the presence of octet $SU(3)$ breaking. Violations of this equal-spacing rule arise at leading order from loop graphs involving the pseudo-Goldstone bosons and behave as $m_s^{3/2}$. As the violation must transform as a **27** under $SU(3)$ we find that the combination of meson masses that enters is numerically very small and we expect the equal-spacing rule to be well satisfied. There is also an equal-spacing rule for the hyperfine mass splittings between the charmed baryons in the **6** and the **6***. This also receives corrections from meson loop graphs but of the form $m_s \log m_s$ (a consequence of heavy quark symmetry). Despite the contribution to individual hyperfine splittings being divergent and requiring the presence of an unknown counterterm the equal spacing rule that holds in the presence of octet $SU(3)$ breaking receives only a finite and numerically small correction. Therefore, we expect that the equal spacing rules that hold in the presence of octet $SU(3)$ breaking , $M_{\Sigma_c} + M_{\Omega_c} - 2M_{\Xi_c} = 0$ and $\delta_{\Sigma_c} + \delta_{\Omega_c} - 2\delta_{\Xi_c} = 0$, are well satisfied in nature. This is not because the loop corrections to the baryons masses are small in chiral perturbation theory (they are not) but because the group structure forces any violation of relations that hold in the presence of octet $SU(3)$ breaking to be small.

The equal-spacing rules and their loop corrections in the charmed baryon sector have direct analogues in the b-baryon sector. We have that $M_{\Sigma_b} + M_{\Omega_b} - 2M_{\Xi_b} = 0$ and $\delta_{\Sigma_b} + \delta_{\Omega_b} - 2\delta_{\Xi_b} = 0$, with only small corrections from meson loops.

Acknowledgements

I would like to thank Jonathan Rosner for stimulating this work and for invaluable discussions and Elizabeth Jenkins for an important discussion. I would also like to thank the Institute for Nuclear theory at the University of Washington for kind hospitality during the course of this work.

REFERENCES

[1] Particle Data Group, L. Montanet et al., Phys. Rev. D50, 1174 (1994); 1995 off year partial update for 1996 edition.

[2] CLEO Collaboration, P. Avery et al., CLNS 95/1352, CLEO 95-14, hep-ex/9508010.

[3] WA89 Collaboration, R. Werding et al., Presented at 27th International Conference on High Energy Physics, Glasgow, July 1994; E. Chudakov et al., talk presented at HQ94, "Heavy Quarks at Fixed Targets", Virginia, October 1994.

[4] SKAT bubble chamber collaboration, V.V. Ammosov et al., presented at International Symposium on Lepton Photon Interactions, Ithaca, 1993.

[5] A. Sakharov, Zh. Eksp. Teor. Fiz. 78, (1980) [Sov. Phys. -JETP 51, 1059 (1980); S. Gasiorowicz and J. L. Rosner, Am. J. Phys. 49, 954 (1981); L.-H. Chan, Phys. Rev. D15, 2478 (1977); L.A. Copley, N. Isgur and G. Karl, Phys. Rev. D20, 768 (1979); K. Maltman and N. Isgur, Phys. Rev D22, 1701 (1980); J.-M. Richard and P. Taxil, Phys. lett. B128, 453 (1983); W.-Y. P. Hwang and D. Lichtenberg, Phys. Rev. D35, 3526 (1986); J.G. Korner and H.W. Siebert, Ann. Rev. Nucl. Part. Sci. 41, 511 (1991); P. Geiger and R.E. Cutkosky, Phys. rev. D 48, 1315 (1993); R. Roncaglia et al., Phys. Rev. D51, 1248 (1995); R. Roncaglia, D.B. Litchtenberg and E. Predazzi, Indiana University Preprint, IUHET-293, January 1995; J.L. Rosner, Enrico Fermi Institute preprint EFI 95-02, Jan. 1995, to appear in Comments on Nuclear and Particle Physics.

[6] J.L. Rosner, Enrico Fermi Institute preprint, EFI-95-48 July 1995.

[7] M.B. Wise Phys. Rev. D45, 2188 (1992); G. Burdman and J. Donoghue, Phys. Letts. B208, 287 (1992); T.M. Yan et al., Phys. Rev. D46, 1148 (1992).

[8] P. Cho, Phys. Letts. B285, 145 (1992).

[9] Z. Guralnik, M. Luke and A.V. Manohar, Nucl. Phys. B390, 474 (1993)

[10] E. Jenkins, Phys. Letts. B315, 431 (1993).

[11] H.-Y. Cheng et al., Phys. Rev. D49, 5857 (1994).

[12] E. Jenkins, Nuc. Phys. B368, 190 (1992).

[13] E. Jenkins, A.V. Manohar and M.B Wise, Caltech Preprint, CALT-68-2000, June 1995, hep-ph/9506356.

[14] E. Jenkins, work in preparation.

[15] M.J. Savage, Phys. Lett. B325, 488 (1994); Phys. Lett. B326, 303 (1994); Phys. Lett. B343, 301 (1995);

[16] E. Jenkins, Nuc. Phys. B412, 181 (1994).